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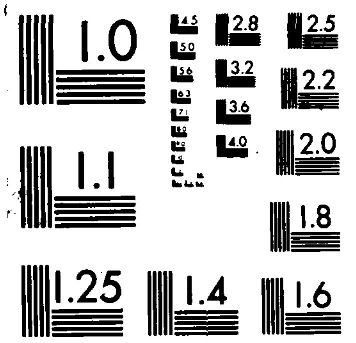
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Small Signal Analysis of the Induced Resonance Electron Cyclotron Maser

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*Science Applications, Intl. Corp., McLean, VA

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The induced resonance electron cyclotron (IREC) maser operates at Doppler upshifted frequencies $\omega = \gamma^2\Omega_o$, with γ the relativistic factor and $\Omega_o = eB_o/\gamma mc$. In addition, the mechanism is insensitive to the beam thermal spreads when the index of refraction $n = k_z c/\omega$ is properly adjusted. A set of simplified equations, describing the electron dynamics near the axis of the configuration, is derived in the limit of small Larmor radius. The linear efficiency is derived for radiation having a Gaussian profile and propagating at an angle α relative to a constant external magnetic field. The start-up current for operation in the submillimeter regime is well within the capabilities of today's long pulse relativistic electron beams.

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CONTENTS

SMALL SIGNAL ANALYSIS OF THE INDUCED RESONANCE ELECTRON CYCLOTRON MASER

I. Introduction

Generation of intense radiation in the microwave regime utilizing electron cyclotron interaction has been proposed independently by a number of researchers in the late 1950's.[1-4] Electrons gyrating in resonance with the radiation field can experience a bunching in the relative wave-particle phase through the dependence of the cyclotron frequency on the relativistic mass. High amplification of the radiation field, known as masing action, results for Doppler shifted frequencies slightly above the electron cyclotron frequency. Electron cyclotron masers, also called gyrotrons,[5-30] have demonstrated efficient high power generation of electromagnetic waves at centimeter wavelengths.

For many purposes it is of practical interest to develop high power generation capability at millimeter and submillimeter wavelengths. Potential areas of application include advanced accelerators, short wavelength radar, plasma heating in fusion reactors and spectroscopy. The shortest wavelength for single mode operation in a closed resonator is tied to the transverse dimension of the cavity. For radiation wavelengths much shorter than the transverse dimensions, a multimode excitation will result from the small frequency separation among nearby modes. The mode selectivity is greatly improved by the use of an open resonator configuration, the quasi-optical maser.[19,20]

A new configuration has recently been proposed[29,30] which utilizes the benefits of the open resonators and at the same time minimizes the detrimental effects of the injected electron beam energy spread. The operating frequency in the induced resonance electron

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cyclotron (IREC) quasi-optical maser is upshifted by a factor γ^2 relative to the relativistic electron cyclotron frequency. It has been shown that for operation at the optimum index of refraction the efficiency is relatively insensitive to the beam energy spread and the sensitivity to the effect of pitch angle spread can be minimized. The index of refraction is adjustable by varying the angle between the resonators (see Fig. (1)) and the guide field, and can be chosen to minimize the effects of finite beam quality. Finally, by spatially tapering the magnetic field the operating efficiency can be increased.

In this paper we limit ourselves to analyzing the small signal efficiency characteristics of such a device. We include the effects of the Gaussian profile for the radiation envelope considering a uniform magnetic field for simplicity. Nonlinear effects and the role of the magnetic field tapering are treated elsewhere.[30]

The remainder of this paper is organized as follows. In Sec. II we describe the field configuration and the equations of motion. In Sec. III we derive the linear energy, power efficiency and start-up current condition. In Sec. IV numerical results and conclusions are presented.

II. Field and Particle Dynamics

The configuration for the induced resonance electron cyclotron (IREC) quasi-optical maser is shown schematically in Fig. 1. The interaction cavity is formed by two quasi-optical resonators intersecting at an angle 2α where α is the angle relative to the external magnetic field B_o in the z direction.

The beam radius is much smaller than the Gaussian width r_o (spot size) for the radiation envelope. In the limit of small Larmor radius ρ compared to the perpendicular wavelength $k_\perp \rho << 1$ we can approximate the vector potential in the interaction regime by

$$\underset{\sim}{A}_T = A_R(z) \exp[i\Phi(z,t)] \frac{1}{2} (\hat{e}_x + i\hat{e}_y) + cc$$

$$+ A_L(z) \exp[i\Phi(z,t)] \frac{1}{2} (\hat{e}_x - i\hat{e}_y) + cc. \qquad (1)$$

Since we are interested in the synchronous interaction of the gyrating electrons with the radiation, we have kept only the forward propagating wave component $\Phi(z,t) = k_z z - \omega t + \Phi_o$. The amplitudes A_R and A_L for the right- and left-handed polarized wave component, respectively are given by

$$A_{R,L}(z) = A^o_{R,L} \exp[- z^2/L^2],$$

$$A^o_{R,L} = A_o (\cos\alpha \pm 1),$$

$$L = r_o/\sin\alpha, \qquad (2)$$

where A_o and r_o are the amplitude and spot size for each individual resonator beam.

We use the guiding center description for the particle orbits

$$x = x_g + \rho \sin\zeta, \quad y_g = y - \rho \cos\zeta,$$

$$p_x = p_{gx} + p_\perp \cos\zeta, \quad p_y = p_{gy} + p_\perp \sin\zeta, \qquad (3)$$

to obtain the nonlinear relativistic equations of motion. In this representation (x_g, y_g) and (p_{gx}, p_{gy}) denote the transverse coordinates

and momentum of the particle's guiding center, ρ is the Larmor radius, $p_\perp$ is the magnitude of the transverse momentum and ζ is the momentum space angle. We assume that x, y, p_x, p_y, ρ and $p_\perp$ are slowly changing, on the spatial scale of a gyroperiod. An additional condition for ignoring finite $k_\perp$ effects is that the guiding center shift in the x direction be small $k_\perp \Delta x << 1$, valid for $\alpha << 1$ where $k_\perp$ is k sinα. Using the Lorentz force equation together with Maxwell's equations and retaining only the right-hand polarized wave component the nonlinear relativistic equations of motion are cast into the form

$$u_\perp' = -[(\omega\gamma/cu_z)-k_z]a(z)\cos\psi + a'(z)\sin\psi, \tag{4a}$$

$$u_z' = -(u_\perp/u_z)[k_z a(z)\cos\psi + a'(z)\sin\psi], \tag{4b}$$

$$\psi' = -(\gamma\Delta\omega/cu_z) + (1/u_\perp)\left[\left((\omega\gamma/cu_z)-k_z\right)a(z)\sin\psi + a'(z)\cos\psi\right], \tag{4c}$$

The prime (') in Eqs. (4) signifies the d/dz derivative, $\underset{\sim}{u} = \underset{\sim}{p}/m_o c = \gamma\underset{\sim}{v}/c$, $\gamma = (1 + u_\perp^2 + u_z^2)^{1/2}$ is the relativistic mass factor, $a(z) = |e|A_R(z)/m_o c^2$ is the normalized radiation amplitude, $\psi = \zeta + \phi$ is the relative phase between the radiation field and particle, $n = ck_z/\omega = \cos\alpha$ is the refractive index associated with the radiation field, $\Delta\omega = [\omega(1-n\beta_z) - \Omega_o/\gamma]$ is the frequency mismatch term and $\Omega_o = |e|B_o/m_o c$ is the nonrelativistic electron cyclotron frequency. Using Eqs. (4) the rate of change of γ is given by

$$\gamma' = -\omega(u_\perp/cu_z)a(z)\cos\psi. \tag{5}$$

The frequency mismatch $\Delta\omega$ and its dependence on the particle energy through the relativistic correction γ, provide the mechanism for the masing action (phase bunching).

III. Efficiency

One of the central issues concerning maser operation is the efficiency of the configuration. Efficiency calculations have been carried out for various configurations in the general categories of the closed resonator gyrotron[5-18,21-28] or the open resonator quasi-optical maser.[19,20] While it is generally recognized that nonlinear saturation mechanisms are very important for the full power operation, it is useful to carry out the small signal efficiency calculation in order to compute the start-up current. Expressions for the small amplitude efficiency, obtained in closed form, provide some guidelines in selecting the optimum operating parameters.

Assuming steady state operation, with the number of particles crossing the unit area per unit time $n_o v_z$ being constant, the efficiency can be defined by

$$\eta_E = -\left\langle \frac{\gamma_f - \gamma_o}{\gamma_o - 1} \right\rangle = \langle \gamma_o - 1 \rangle^{-1} \int d^3 \underline{p}_o \underline{f}_o(\underline{p}_o) \int_{-\infty}^{\infty} dz \, \frac{\partial \gamma}{\partial z} \, . \qquad (6)$$

In Eq. (6), the bracket < > signifies the average over the initial distribution in phase space, the subscript $\pm\infty$ stands for the initial and final values at $z = \pm\infty$ respectively and $\partial\gamma/\partial z$ is a function of the initial conditions $\gamma = \gamma(z; p_{\perp o}, p_{zo}, \psi_o)$. In the cold beam limit with the initial distribution function given by $f_o(p_\perp, p_z, \psi) = (n_o/2\pi p_\perp)$ $\delta(p_\perp - p_{\perp o})\delta(p_z - p_{zo})$ the average reduces to an average over $\psi_o = \zeta_o + \Phi_o$.

a. Small Signal Efficiency

We proceed to compute the small signal power efficiency by evaluating the right-hand side of Eq. (6) using Eq. (5). A first order expansion for the quantities $u_\perp = u_\perp^{(o)} + u_\perp^{(1)}$, $\gamma = \gamma^{(o)} + \gamma^{(1)}$, $\psi = \psi^{(o)} + \psi^{(1)}$ will suffice for a quadratic expression in the wave amplitude a. The integrand on the right-hand side of Eq. (6) is expanded using the linear solutions from Eqs. (4a)-(4c). The evaluation of the final result is considerably simplified by performing the phase space average over the angle ψ_o before the spatial integration over z. Expanding the products of the trigonometric terms inside the integral in Eq. (6) into sums and averaging over ψ_o leads to

$$\left\langle \int_{-\infty}^{\infty} dz \frac{\partial\gamma}{\partial z} \right\rangle = -\left(\frac{\omega}{cu_{zo}}\right)\left[\left(1 + \frac{1}{2}\frac{u_{\perp o}^2}{u_{zo}^2}\right)\int_{-\infty}^{\infty} dz a(z)\int_{-\infty}^{z} dz' \frac{da(z')}{dz'} \sin \Delta_o(z-z')\right.$$

$$+ \left\{k_z\left(1 + \frac{1}{2}\frac{u_{\perp o}^2}{u_{zo}^2}\right) - \frac{\omega\gamma_o}{cu_{zo}}\right\}\int_{-\infty}^{\infty} dz a(z)\int_{-\infty}^{\infty} dz' \; a(z')\cos\Delta \; (z-z')$$

$$- \frac{1}{2}\left\{\frac{\omega^2}{c^2}\frac{u_{\perp o}^2}{u_{zo}^2} - k_z\frac{u_{\perp o}^2}{u_{zo}^2}\left(\frac{\omega\gamma_o}{cu_{zo}} - \frac{\Omega_o}{cu_{zo}}\right)\right\}\int_{-\infty}^{\infty} dz a(z)\int_{-\infty}^{z} dz'\int_{-\infty}^{z'} dz'' a(z'')\sin\Delta_o(z-z'')$$

$$\left. + \frac{1}{2}\frac{u_{\perp o}^2}{u_{zo}^2}\left(\frac{\omega\gamma_o}{cu_{zo}} - \frac{\Omega_o}{cu_{zo}}\right)\int_{-\infty}^{\infty} dz a(z)\int_{-\infty}^{z} dz'\int_{-\infty}^{z'} dz'' \frac{da(z'')}{dz''} \cos\Delta_o(z-z'')\right], \qquad (7)$$

where $\Delta_o = \frac{\Omega_o}{cu_{zo}} - \left(\frac{\omega\gamma_o}{cu_{zo}} - k_z\right) = -\frac{\Delta\omega_o}{v_{zo}}$, $a(z) = a_o\exp[-z^2/L^2]$

and $a_o = |e| \, A_R^o/m_oc^2$.

We evaluate the remaining integrals in Eq. (14) and express the final result in terms of the parameters $\xi = \omega\tau = (\omega\gamma_o/cu_{zo})L$, τ being the transit time through the interaction regime, and the relative frequency mismatch $\Delta\omega_o/\omega$. We find

$$\eta_p = \frac{\pi}{2}\,\frac{a_o^2\,\xi^2}{\gamma_o(\gamma_o-1)}\left\{\left(1+\theta_o^2\right)n\beta_{zo}-1\right.$$

$$+\left[\frac{1}{2}\,\xi^2\beta_{\perp o}^{\,2}(1-n^2)+\left(1+\theta_o^2\right)\right]\frac{\Delta\omega_o}{\omega}-\theta_o^2\,n\beta_{zo}\,\xi^2\left(\frac{\Delta\omega_o}{\omega}\right)^2 \tag{8}$$

$$\left.-\frac{\theta_o^2}{2}\;\xi^2\left(\frac{\Delta\omega_o}{\omega}\right)^3\right\}e^{-\frac{1}{2}\,\xi^2\,\frac{\Delta\omega_o^2}{\omega^2}},$$

with $\theta_o = u_{\perp o}/u_{zo}$, the initial pitch angle.

The efficiency is proportional to $\exp[-1/2\ \xi^2\ \Delta\omega_o^2/\omega^2]$ where exponent $\xi(\Delta\omega_o/\omega)$ is equal to $\Delta\omega_o\tau$, the advance in the relative phase $\Delta\psi_o$ between the wave and the particle over the interaction regime. For typical values of $\xi \gg 1$ and $\Delta\omega_o/\omega \ll 1$ the expression in braces in Eq. (8) is simplified to

$$\left\{\cdots\right\} \simeq \left(1+\theta_o^2\right)n\beta_{zo}-1+\frac{\xi_o^2\beta_{\perp o}^{\,2}}{2}\,\frac{\Delta\omega_o}{\omega}-\theta_o^2\beta_{zo}\,\frac{\xi_o^2}{\sin^2\alpha}\left(\frac{\Delta\omega_o}{\omega}\right)^2, \tag{9}$$

where $\xi_o^{\,2} = \xi^2\,(1-n^2)$ is independent of α. In (9) we have omitted the small terms that originate from the gradient terms $\partial a/\partial z$ in the equations of motion. Treating (9) as a quadratic form in $\Delta\omega/\omega$ we find the regime for positive efficiency, given by

$$2\left(1-\left(1+\theta_o^2\right)n\beta_{zo}\right)\left(\beta_{\perp o}^2\xi_o^2\right)^{-1} < \frac{\Delta\omega_o}{\omega} < \beta_{\perp o}^{\,2}(1-n^2)\left(n\ \theta_o^2\beta_{zo}\right)^{-1}. \tag{10}$$

The upper limit in $\Delta\omega_o/\omega$ is due to a finite n and results from the negative contribution of the quadratic term $(\Delta\omega/\omega)^2$ that overtakes the positive contribution of the linear term $\Delta\omega/\omega$ for small angles $\sin^2\alpha < (2n\theta_o^2\beta_{zo}/\beta_{\perp o}^2)\,(\Delta\omega_o/\omega)$.

In order to determine the maximum efficiency within the positive regime, we parameterize Eq.(8) as a function of $x = \xi\,\Delta\omega/\omega$, since the exponential is the main factor limiting efficiency. Setting $d\eta/dx = 0$, we obtain

$$c_3x^3 - c_2\,x^2 - c_1x + c_o = 0\ , \tag{11}$$

with $c_1 = (1 + 3\theta_o^2)\beta_{zo}\cos\alpha - 1$, $c_3 = \theta_o^2\beta_{zo}\cos\alpha$ and $c_2 = c_o = (1/2)\,\beta_{\perp o}^2\xi_o\sin\alpha$. Observing that the terms proportional to c_1 and c_3 can be omitted provided that $c_o = c_2 >> c_3 \sim c_1 \sim 1$ or

$$\sin\alpha >> \frac{\theta_o^2\beta_{zo}}{\beta_{\perp o}^2\xi_o}, \tag{12}$$

we can show that $x \simeq 1$. In the special case $x = 1$, we obtain the maximum efficiency

$$\eta_{max} = \frac{\pi}{4}\,a_o^2e^{-1/2}\,\frac{\beta_{\perp o}^2\xi_o^3}{\gamma_o(\gamma_o-1)}\,(\sin\alpha)^{-1}. \tag{13}$$

The overall efficiency increases with decreasing α (increasing index of refraction) provided that inequality Eq. (12) remains valid. For very small α Eq. (13) fails and a solution of the cubic Eq. (11) is necessary.

b. Start-up Current

We are in position now to calculate the start-up beam current utilizing the power efficiency coefficient. Amplification of the electromagnetic field energy will result if

$$\eta P_b > \frac{d\varepsilon}{dt}, \tag{14}$$

where ε is the total electromagnetic energy stored in both cavities $\varepsilon = \int U_R dV = 2V(\omega^2/c^2)(A_o^2/4\pi)$, $V=\pi r_o^2 L_T$, $d\varepsilon/dt = (\omega/Q)\varepsilon$, Q is the quality factor for the cavity and P_b is the electron beam power.

The optimum power efficiency η_{max} is given by Eq. (13). The cavity Q is given by

$$Q = \frac{2\pi}{1-R_{ef}} \frac{L_T}{\lambda}, \tag{15}$$

where L_T is the effective resonator length and λ the wavelength. Combining Eqs. (13), (14), (15) and expressing A_o in terms of a_o from Eq. (2) we obtain

$$P_b > \frac{\lambda}{r_o} (1-R_{ef}) \frac{\exp(\frac{1}{2})}{4\pi^2} \frac{m_o^2 c^5}{|e|^2} \frac{\beta_{zo}^3 \gamma_o(\gamma_o-1)}{\beta_{\perp o}^2} \frac{2\sin\alpha}{(1+\cos\alpha)^2}, \tag{16}$$

where $P_b = I_b V_b$, I_b is the current and V_b is voltage of the electron beam. For typical parameters $V_b = 0.25 \times 10^6$eV, $\lambda/r_o = 10^{-1}$, $1-R_{ef} = 0.1$, $\gamma_o = 1.5$, $\beta_{zo} \simeq 0.64$, $\beta_{\perp o} \simeq (\sqrt{3}\gamma_o)^{-1}$, and the optimum operation angle $\alpha \cong 45°$, the start-up current is

$$I_b \gtrsim 4.6 \text{ A}.$$

IV. Conclusion and Summary

We have performed the small signal analysis for an oscillator configuration capable of generating radiation in the millimeter and the submillimeter regime. The threshold for the start-up current was found to be well within the existing capabilities of today's long pulse mildly relativistic beams. Our theoretical linear efficiency results are plotted as solid lines in Figs. 2-4 against the numerical results (dots) obtained by direct integration of the fully nonlinear Eqs. (4) for small wave amplitude. Plots of the linear efficiency as a function of the controlling parameter $\xi\, \Delta\omega/\omega$ for constant radiation amplitude a_o and constant spot size r_o are shown in Fig. 2, with each curve corresponding to a different index of refraction $n = \cos\alpha$. The maximum efficiency for all plots occurs at $\xi\, \Delta\omega/\omega \simeq 1$ in agreement with Eq. (13). Small signal efficiency increases with increasing $n = \cos\alpha$ roughly proportionally to the length of the interaction regime $L = r_o/\sin\alpha$. In Fig. 3, the optimum index of refraction[29-30] $n = \beta_{zo}/(1-\beta_{\perp o}^2)$, to minimize the effects of beam energy spread, is held constant, and the interaction length L is changed by increasing the width of the radiation envelope r_o. Figure 4 is a comparison of the theoretical small signal efficiency with the numerically calculated nonlinear efficiency as a function of wave amplitude a_o. The agreement is good for $a_o \leq 3 \times 10^{-4}$. Nonlinear saturation occurs for $a_o > 1 \times 10^{-3}$. Obtaining the scaling of the efficiency in the nonlinear regime is not possible analytically. Numerical studies of the high power performance, however, have demonstrated good nonlinear efficiency.

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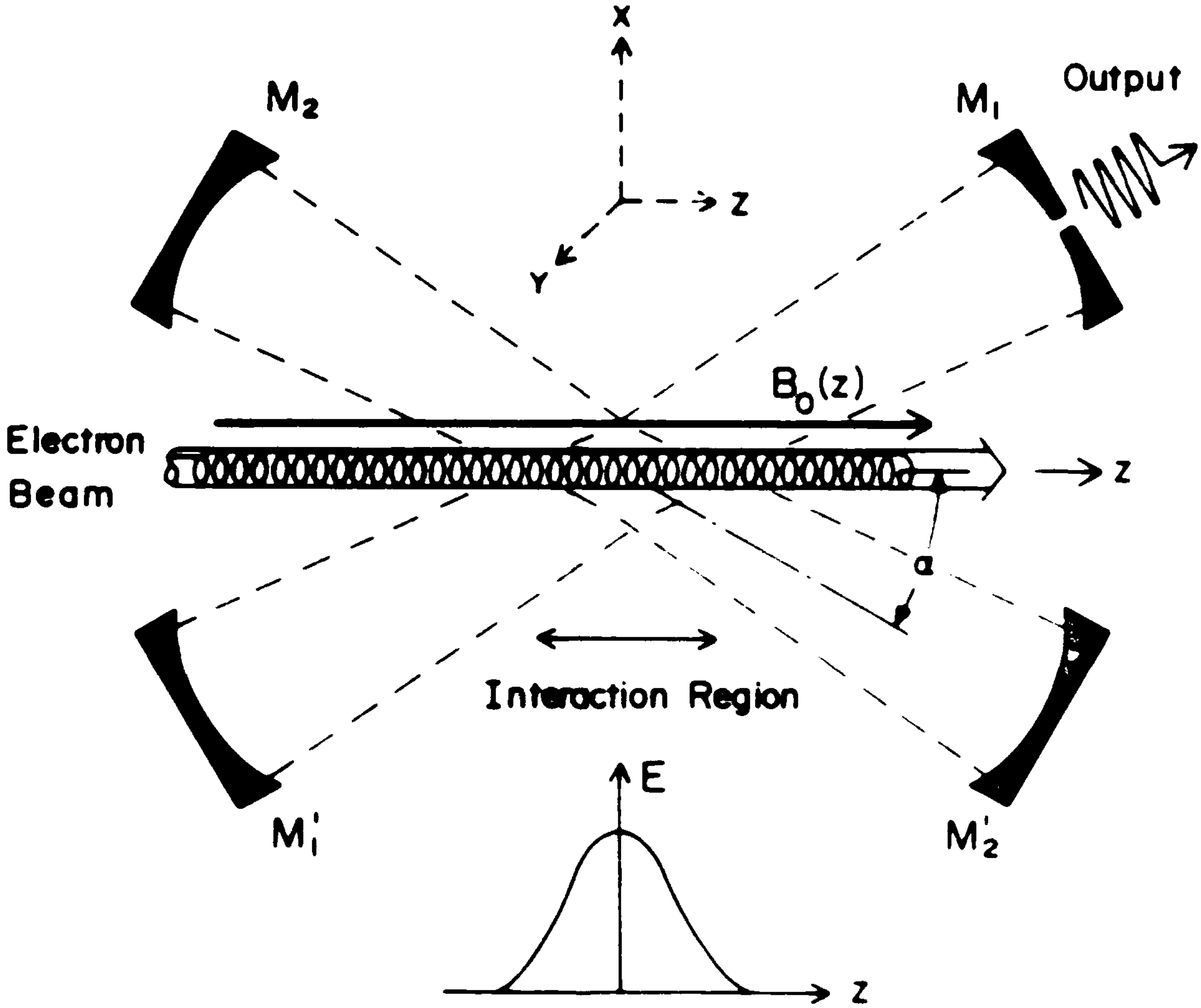

Figure 1. The configuration of the Induced Electron Resonance Maser.

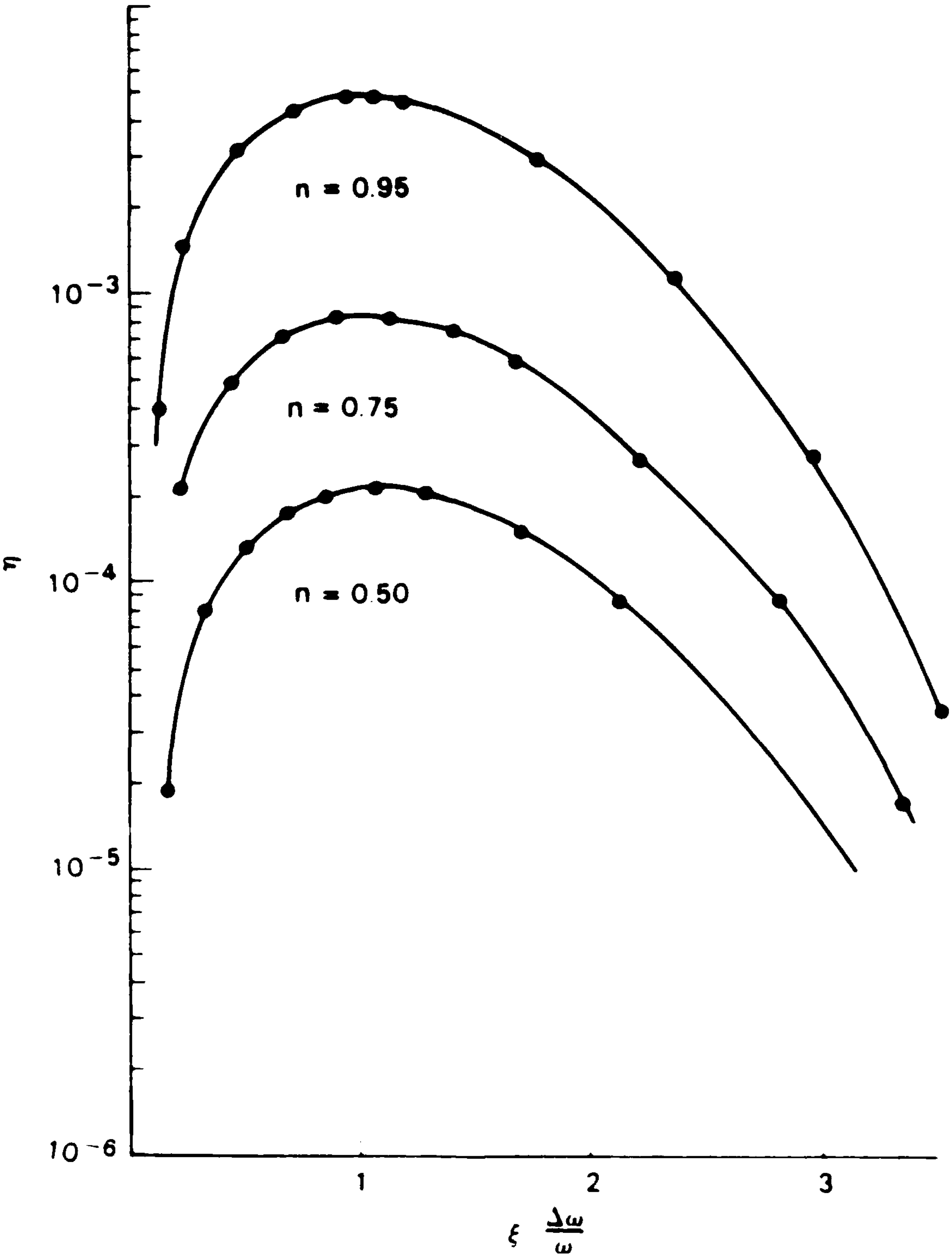

Figure 2. Theoretical (solid line) and numerical (dots) plots of linear efficiency η versus $\xi\,\Delta\omega/\omega$ for various values of index of refraction n with constant amplitude $a_0 = 5\times10^{-5}$ and $\gamma = 1.5$.

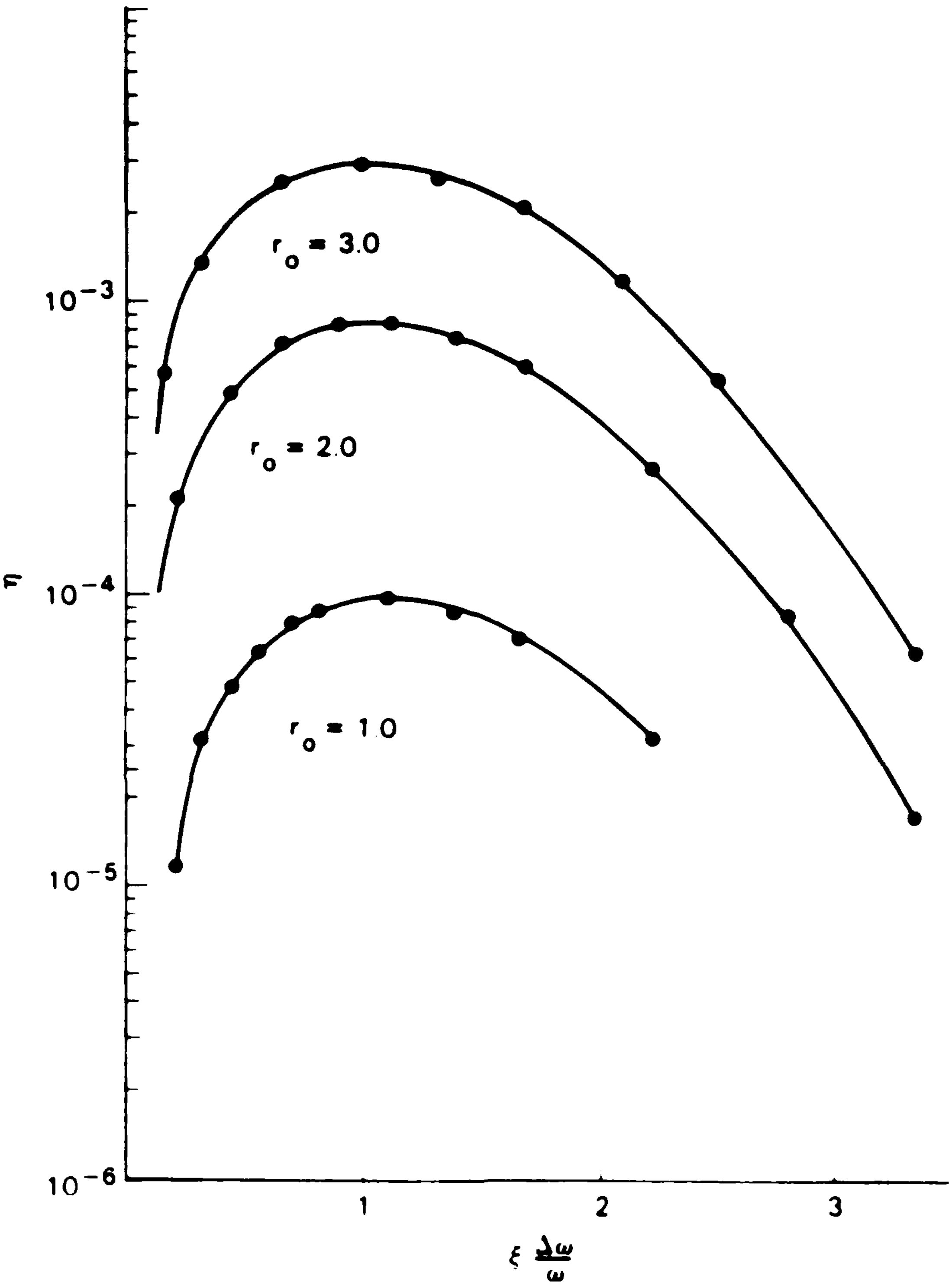

Figure 3. Theoretical (solid line) and numerical (dots) plots of linear efficiency η versus $\xi\,\Delta\omega/\omega$ for various Gaussian widths r_0 with constant refraction index n_{opt} and $a_0 = 5\times10^{-5}$, $\nu = 1.5$.

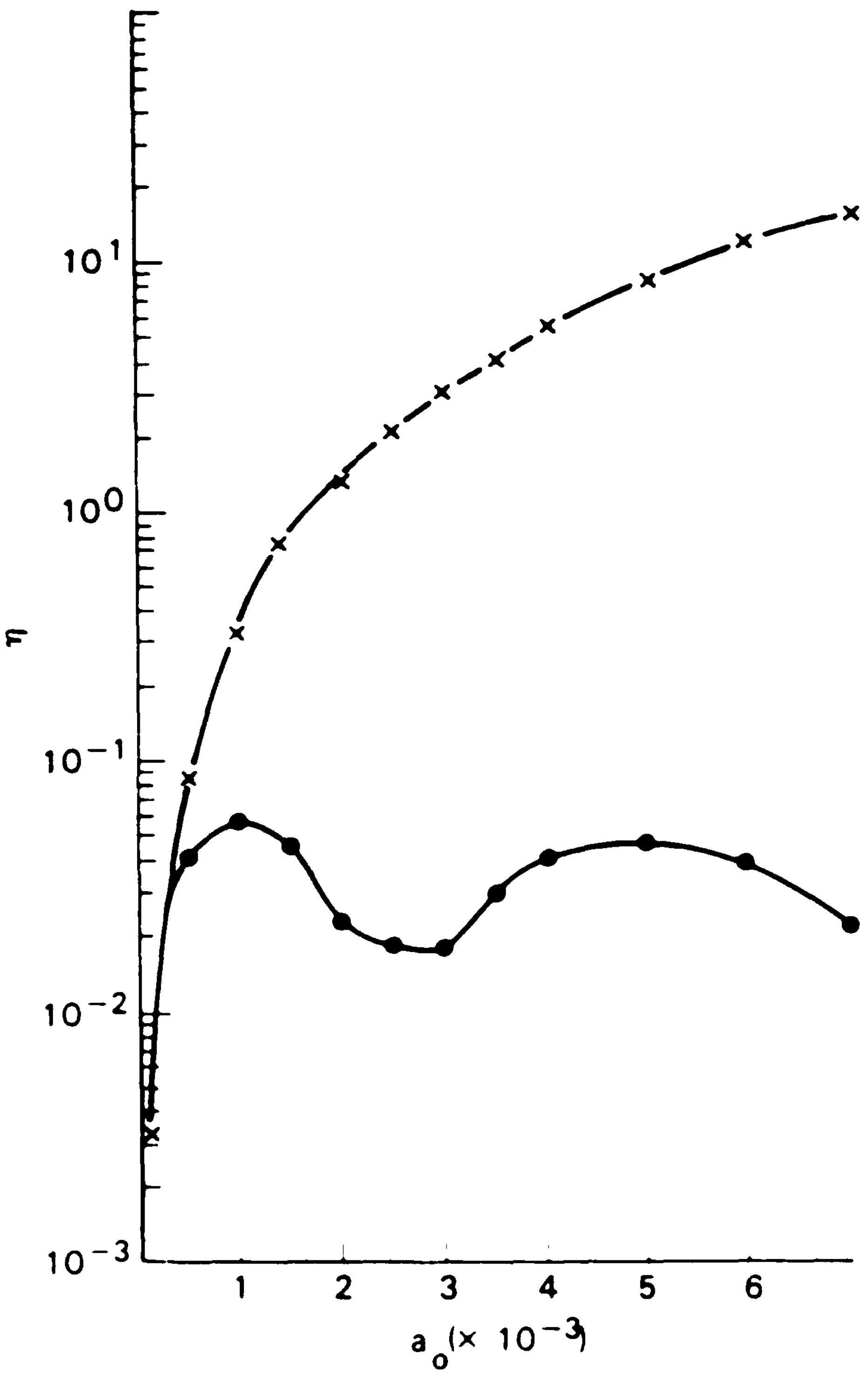

Figure 4. Comparison of linear (crosses) versus nonlinear (dots) efficiency η as a function of a_o for $\xi\,\Delta\omega/\omega = 1$, $n = 0.75$ and $\gamma = 1.5$.

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